

125 GeV Higgs boson mass and muon $g - 2$ in 5D MSSMNobuchika Okada^{1,*} and Hieu Minh Tran^{1,2,†}¹*Department of Physics and Astronomy, University of Alabama, Tuscaloosa, Alabama 35487, USA*²*Hanoi University of Science and Technology, 1 Dai Co Viet Road, Hanoi, Vietnam*

(Received 21 July 2016; published 26 October 2016)

In the minimal supersymmetric standard model (MSSM), the tension between the observed Higgs boson mass and the experimental result of the muon $g - 2$ measurement requires a large mass splitting between stops and smuons/charginos/neutralinos. We consider a five-dimensional (5D) framework of the MSSM with the Randall-Sundrum warped background metric, and show that such a mass hierarchy is naturally achieved in terms of geometry. In our setup, the supersymmetry is broken at the ultraviolet (UV) brane, while all the MSSM multiplets reside in the 5D bulk. An appropriate choice of the bulk mass parameters for the MSSM matter multiplets can naturally realize the sparticle mass hierarchy desired to resolve the tension. The gravitino is localized at the UV brane and hence becomes very heavy, while the gauginos spreading over the bulk acquire their masses suppressed by the fifth dimensional volume. As a result, the lightest sparticle neutralino is a candidate for the dark matter as usual in the MSSM. In addition to reproducing the SM-like Higgs boson mass of around 125 GeV and the measured value of the muon $g - 2$, we consider a variety of phenomenological constraints, and present the benchmark particle mass spectra that can be explored at the LHC Run-2 in the near future.

DOI: 10.1103/PhysRevD.94.075016

I. INTRODUCTION

The standard model (SM) of elementary particles has been tested to a very high accuracy. It can be considered to be completed in the sense that all particles in the SM have been observed and their properties have been confirmed to be consistent with the SM expectations. However, there are still problems that this model itself cannot address, and hence that motivate us to go beyond the SM. The discovery of the Higgs boson not only fulfils the SM particle content, but also provides us with a hint that the SM needs to be extended according to the requirement of naturalness. Cosmological observations confirmed the existence of dark matter that goes beyond the SM prediction. On the other hand, the measurement for the muon anomalous magnetic dipole moment ($g - 2$) [1] reveals a 3–4 σ discrepancy between the measured central value and the SM prediction [2].

Supersymmetry (SUSY) has been investigated for a long time as one of the most promising candidates beyond the SM. The gauge hierarchy problem, in other words, the instability of the electroweak scale under quantum corrections, can be solved by a SUSY extension of the SM such as the minimal supersymmetric SM (MSSM). Assuming the conservation of R parity, SUSY models can provide good candidates for the cold dark matter in the Universe. Additionally, the muon $g - 2$ receives contributions from superpartners

(smuons, charginos, and neutralinos) and shifts to the allowed measured interval when their masses lie at the electroweak scale [3].

The Higgs boson mass measurement at the Large Hadron Collider (LHC), together with other experiments, has put severe constraints on SUSY breaking parameters. At the tree level, the SM-like Higgs boson mass m_h is just about the Z boson mass m_Z . To reproduce the Higgs boson mass of about 125 GeV [4], quantum corrections to the Higgs boson mass play a crucial role. The approximate formula for m_h with radiative corrections (in case A_t this is relatively large compared to μ , and $m_A \gg m_Z$) is given by [5]

$$m_h^2 \approx m_Z^2 \cos^2 2\beta + \frac{3}{4\pi^2} y_t^2 m_t^2 \sin^2 \beta \left[\log \left(\frac{m_t^2}{m_t^2} \right) + \frac{X_t^2}{m_t^2} - \frac{X_t^4}{12m_t^4} \right], \quad (1)$$

where $X_t = A_t - \mu \cot \beta$ is the stop mixing parameter. In this formula, a large stop mass plays a crucial role to push up the Higgs boson mass from m_Z to the measured value of about 125 GeV. In many SUSY breaking models, heavy stops imply that other sfermions are also heavy, such that squark masses are of $\mathcal{O}(10 \text{ TeV})$, and slepton masses lie around a few TeV [6]. However, heavy smuons make their loop contributions to the muon anomalous magnetic moment, $a_\mu = \frac{1}{2}(g_\mu - 2)$, too small to explain the 3–4 σ discrepancy. This fact can be easily seen in the formula of SUSY contributions to the muon anomalous magnetic moment [7],

*okadan@ua.edu

†hieu.tranminh@hust.edu.vn

$$\Delta a_\mu = \frac{\alpha m_\mu^2 M_2 \tan \beta}{4\pi \sin^2 \theta_W m_{\tilde{\mu}_L}^2} \left[\frac{f_\chi(M_2^2/m_{\tilde{\mu}_L}^2) - f_\chi(\mu^2/m_{\tilde{\mu}_L}^2)}{M_2^2 - \mu^2} \right] + \frac{\alpha m_\mu^2 M_1 \tan \beta}{4\pi \cos^2 \theta_W (m_{\tilde{\mu}_R}^2 - m_{\tilde{\mu}_L}^2)} \times \left[\frac{f_N(M_1^2/m_{\tilde{\mu}_R}^2)}{m_{\tilde{\mu}_R}^2} - \frac{f_N(M_1^2/m_{\tilde{\mu}_L}^2)}{m_{\tilde{\mu}_L}^2} \right], \quad (2)$$

where the loop functions are defined as

$$f_\chi(x) = \frac{x^2 - 4x + 3 + 2 \ln x}{(1-x)^3},$$

$$f_N(x) = \frac{x^2 - 1 - 2x \ln x}{(1-x)^3}. \quad (3)$$

For $x = \mathcal{O}(1)$, they are of order 1, for example, $f_\chi(1) = -\frac{2}{3}$ and $f_N(1) = -\frac{1}{3}$.

To yield $\Delta a_\mu \sim 10^{-9}$ to fill the discrepancy between the experimental result and the SM prediction, light smuons and charginos/neutralinos are necessary, while a large $\tan \beta$ works to enhance Δa_μ . A solution to the tension between the Higgs boson mass and the muon $g-2$ may come from a large mass splitting between stops (squarks in general) and smuons (leptons in general)/charginos/neutralinos [8]. According to Eq. (1), besides heavy stops the Higgs boson mass can also be improved by large X_t [9–11]. There are also other proposals to solve this tension [12].

In this paper, we investigate the MSSM in the five-dimensional (5D) space-time with the Randall-Sundrum (RS) background metric [13]. Originally, the RS model was proposed to solve the gauge hierarchy problem of the SM, where all the SM particles are confined on the so-called infrared (IR) brane at a fixed point of the S^1/Z_2 orbifold on which the fifth dimension is compactified. The large hierarchy between the electroweak and the Planck scales is naturally generated via the so-called warp factor induced by the RS warped background metric. Soon after the original work, the RS model was extended to have the SM fields reside in the bulk while the SM Higgs field is confined on the IR brane to maintain the solution to the gauge hierarchy problem [14]. SUSY extensions of the RS model then came up with the component field formulation [15] and the superfield formulation [16]. This context provides not only an elegant explanation of the diversity of particle masses [17], but also a variety of possibilities for SUSY breaking mediation mechanisms with the fifth dimension.¹ In SUSY RS models, the SUSY can solve the gauge hierarchy problem as usual, so that a very strong warp factor is not necessary. SUSY RS models have the

¹Beside the RS scenario, the flat extradimension scenario also brings up an interesting landscape for discussions on SUSY breaking. See, for example, [18] and references therein.

ability to simultaneously solve both hierarchy problems: the gauge hierarchy and the fermion mass (Yukawa) hierarchy. On the other hand, the AdS/CFT correspondence brings an interesting point to our study setup. This conjecture maps the physics in the AdS₅ space on to its dual four-dimensional (4D) picture with a strongly coupled conformal field theory [19]. A connection of the models to string theories would be possible [20].

Our setup is similar to Ref. [21] where all matter and gauge superfields reside in the bulk. But we arrange the Higgs superfields to reside in the bulk as well, while the SUSY breaking hidden sector is confined on the so-called ultraviolet (UV) brane. In particular, quark superfields localize around the UV brane, while slepton superfields delocalize from the UV brane. We consider that the 5D MSSM can have a universal coupling between the MSSM multiplets and the hidden sector field. However, because of this geometrical configuration of wave functions, we can naturally realize a large mass splitting between squarks and sleptons. Since the gravitino, which is the superpartner of the massless 4D graviton, localizes around the UV brane, its mass is large. Hence, the lightest neutralino serves as a dark matter candidate as usual in the MSSM.

Taking into account a variety of experimental constraints, we find bulk field configurations that are suitable to account for the muon $g-2$ while reproducing the observed SM-like Higgs boson mass.

The structure of this paper is as follows. In Sec. II, we describe the 5D MSSM with bulk superfields in the Randall-Sundrum background metric, and derive the 4D MSSM as a low-energy effective theory. In Sec. III, we consider various phenomenological constraints, and present the benchmark particle spectra of the model. Section IV is devoted to conclusions.

II. 5D MSSM IN THE RS BACKGROUND METRIC

We consider a 5D space-time (x^μ, y) , where the fifth dimension is compactified on the S^1/Z_2 orbifold, and y defined in the range of $-\pi \leq y \leq \pi$ is an angle of S^1 with a radius R . Because of Z_2 parity, y is identified with $-y$, so that the orbifold possesses two fixed points at $y=0$ and $y=\pi$. Introducing two “3-branes” located at these orbifold fixed points and assigning suitable brane tensions to them, a solution to the Einstein’s equation is found to be [13]

$$ds^2 = e^{-2R\sigma} \eta_{\mu\nu} dx^\mu dx^\nu - R^2 dy^2, \quad (4)$$

where $\sigma = k|y|$, and k is the anti-de Sitter (AdS) curvature. The 4D Minkowski space is realized as a slice of this AdS₅ space.

With this metric, a relation between the 5D Planck mass (M_5) and the Planck mass (M_4) in the 4D effective theory is given by

$$M_4^2 = \frac{M_5^3}{k}(1 - \omega^2) \approx \frac{M_5^3}{k}, \quad (5)$$

where $\omega = e^{-kR\pi}$ is the so-called warp factor, and we have assumed $\omega \ll 1$. In the following calculation, we simply take $k \approx M_5 \approx M_4$. Because of the warped metric, an effective cutoff on the UV brane at $y = 0$ is M_5 itself, while that on the IR brane at $y = \pi$ is warped down to $M_{\text{cut}} = M_5\omega$. This effective cutoff at the IR brane is used to solve the gauge hierarchy problem in the original paper by Randall and Sundrum [13].

In this paper, we assume that all the MSSM superfields propagate in the 5D bulk, while the SUSY breaking hidden sector is confined on the UV brane. Since the gravity multiplet localizes toward the UV brane, the gravitino interacts strongly with the hidden sector field due to the large overlapping between their wave functions. Therefore, when the SUSY is broken, the gravitino acquires a large mass. On the other hand, the MSSM gauge multiplets have flat configurations in the fifth dimensions, and gaugino masses are suppressed by the volume of the fifth dimension. Hence the lightest neutralino is a dark matter candidate as usual in the MSSM. The wave function configuration for the MSSM matter multiplets depends on the bulk mass parameters. An sfermion mass becomes smaller, as its wave function is more delocalized from the UV brane.

Now we describe our $N = 1$ 5D MSSM Lagrangian in terms of familiar ‘‘4D-like’’ $N = 1$ supermultiplets.² In the Kaluza-Klein decomposition, only the Z_2 -even 5D fields have massless modes. After integrating over the extra fifth dimension y , these massless zero modes are identified as usual 4D MSSM supermultiplets.

The 5D action for a bulk vector multiplet is described by using a 4D-like vector superfield

$$V(x, y, \theta) = -\theta\sigma^\mu\bar{\theta}A_\mu(x, y) - i\bar{\theta}^2\theta\lambda_1(x, y) + i\theta^2\bar{\theta}\bar{\lambda}_1(x, y) + \frac{1}{2}\bar{\theta}^2\theta^2 D(x, y), \quad (6)$$

and a 4D-like adjoint chiral superfield

$$\chi(x, y, \theta) = \frac{1}{\sqrt{2}}(\Sigma(x, y) + iA_5(x, y)) + \sqrt{2}\theta\lambda_2(x, y) + \theta^2 F_\chi(x, y). \quad (7)$$

Under the Z_2 parity, the former is even while the latter is odd. The gauge invariant 5D action for the gauge multiplet is given by

²In this paper, we assume some mechanism to stabilize the fifth dimensional radius, and simply replace the radion supermultiplet as the fifth dimensional radius R . We refer the reader to Ref. [22] for a very simple mechanism to stabilize the radion potential.

$$\begin{aligned} S_5^{\text{gauge}} = & \int d^4x \int_{-\pi}^{\pi} dy \left\{ \frac{1}{4g_5^2} \int d^2\theta R \text{Tr}[W^\alpha W_\alpha] + \text{H.c.} \right. \\ & + \frac{2}{g_5^2} \int d^4\theta \frac{e^{-2R\sigma}}{2R} \text{Tr} \left[\{e^{V/2}, \partial_y e^{-V/2}\} \right. \\ & \left. \left. + \frac{1}{\sqrt{2}} (e^{V/2} \chi^\dagger e^{-V/2} + (e^{-V/2} \chi e^{V/2}))^2 \right\} \right], \quad (8) \end{aligned}$$

where the 5D gauge coupling g_5 has the mass dimension of $-1/2$. Since the wave function of the vector superfield zero mode is found to be independent of the y -coordinate, we have rescaled it as $V \rightarrow V/\sqrt{2\pi R}$, by which the zero mode is canonically normalized in the 4D effective theory with the relation between the 5D and 4D gauge couplings, $g_5 = \sqrt{2\pi R}g_4$.

A hypermultiplet in the bulk is used to describe matter and Higgs multiplets, which is decomposed into a pair of vectorlike chiral superfields Φ and Φ^c ,

$$\Phi(x, y, \theta) = \phi(x, y) + \sqrt{2}\theta\psi(x, y) + \theta^2 F_\Phi(x, y), \quad (9)$$

$$\Phi^c(x, y, \theta) = \phi^c(x, y) + \sqrt{2}\theta\psi^c(x, y) + \theta^2 F_{\Phi^c}(x, y). \quad (10)$$

Under the Z_2 parity, we assign an even parity for Φ while Φ^c is odd. The 5D action for the hypermultiplet is given by

$$\begin{aligned} S_5^{\text{matter}} = & \int d^4x \int_{-\pi}^{\pi} dy \left\{ \int d^2\theta R e^{-2R\sigma} (\Phi^\dagger e^{-V} \Phi + \Phi^c e^V \Phi^{c\dagger}) \right. \\ & + \int d^2\theta e^{-3R\sigma} \Phi^c \left[\partial_y - \frac{1}{\sqrt{2}} \chi - \left(\frac{3}{2} - c_\Phi \right) R\sigma' \right] \Phi \\ & \left. + \text{H.c.} \right\}, \quad (11) \end{aligned}$$

where c_Φ is a bulk mass parameter. Because of the Z_2 parity assignment of Φ and Φ^c , only the Z_2 -even chiral multiplet Φ has a zero mode in the Kaluza-Klein decomposition. By solving the SUSY vacuum condition,

$$\left[\partial_y - \left(\frac{3}{2} - c_\Phi \right) R\sigma' \right] \Phi = 0, \quad (12)$$

we find the zero-mode wave function as

$$\Phi(x, y, \theta)|_{\text{zero-mode}} = \hat{\Phi}(x, \theta) e^{\left(\frac{3}{2} - c_\Phi\right) R\sigma}. \quad (13)$$

Here, the 4D chiral superfield $\hat{\Phi}(x, \theta)$ has a mass dimension $\frac{3}{2}$. The canonically normalized chiral superfield $\varphi_0(x, \theta)$ in the 4D effective theory is given by

$$\hat{\Phi}(x, \theta) = \sqrt{k} C_\Phi \varphi_0(x, \theta), \quad (14)$$

where φ_0 represents the 4D MSSM chiral superfields ($H_u, H_d, Q_i, U_i, D_i, L_i, E_i$), and

TABLE I. R -charge assignments for the 5D MSSM vector multiplet ($a = 1, 2, 3$), hypermultiplets, and the hidden sector field X (R charge of θ is 1). Here, for example, Q_i^h is a Z_2 -even component of the bulk quark doublet hypermultiplet, whose zero mode is identified as the quark doublet chiral superfield in the 4D MSSM. The generation index is denoted as $i = 1, 2, 3$.

Bulk field	V_a	Q_i^h	U_i^h	D_i^h	L_i^h	E_i^h	H_u^h	H_d^h	X
R charge	0	1	1	1	1	1	0	0	0

$$C_\Phi = \sqrt{\frac{(1 - 2c_\Phi)}{2(\omega^{(-1+2c_\Phi)} - 1)}}. \quad (15)$$

Note that the bulk mass parameter c_Φ controls the configuration of the zero mode: for $c_\Phi > 1/2$ ($c_\Phi < 1/2$), the zero mode is localized toward the UV (IR) brane.³

Now we introduce interaction terms among the bulk multiplets and a chiral multiplet in the hidden sector on the UV brane. Because of the 5D $N = 1$ SUSY, such interaction terms can be written only at the orbifold fixed points. In order to forbid phenomenologically dangerous terms such as R -parity violating terms, we introduce an R symmetry with the charge assignments listed in Table I. Here, a chiral superfield X in the hidden sector has been introduced, and we assume that both the SUSY and the R symmetry are broken by a vacuum expectation value (VEV) of the F -component of X , $\langle F_X \rangle \neq 0$.

In the 5D MSSM, a Yukawa coupling is symbolically given by

$$S_5^{\text{Yukawa}} = \int d^4x \int_{-\pi}^{\pi} dy \int d^2\theta e^{-3R\sigma} \frac{1}{M_5^{3/2}} \Phi_1 \Phi_2 \Phi_3 \times [Y_0 \delta(y) + Y_\pi \{\delta(y + \pi) + \delta(y - \pi)\}], \quad (16)$$

where Y_0 and Y_π are dimensionless coupling constants, and $M_5^{3/2}$ is introduced to yield the correct mass dimension. Here, Φ_1 stands for the MSSM Higgs doublets, and the other two Φ_2 and Φ_3 stand for the MSSM matter multiplets. After the y -integration, a 4D effective Yukawa coupling is obtained as

$$Y_4 \approx [Y_0 + Y_\pi \omega^{-(\frac{3}{2} - c_{\Phi_1} - c_{\Phi_2} - c_{\Phi_3})}] C_{\Phi_1} C_{\Phi_2} C_{\Phi_3}, \quad (17)$$

where we have used $k/M_5 \approx 1$. Note that an appropriate choice of the bulk mass parameters can derive an exponentially suppressed Yukawa coupling even for $Y_0, Y_\pi = \mathcal{O}(1)$. Although this feature implies a possibility to naturally explain the Yukawa hierarchy in the SM, in this paper we

³In order to have the canonical Kahler potential in Eq. (11), we have redefined the hypermultiplet in Eq. (13) as $\Phi \rightarrow \Phi e^{-R\sigma}$ when discussing field localization in the extra dimension.

do not attempt to explain the Yukawa hierarchy, but concentrate on soft SUSY breaking parameters.

Let us consider R -symmetric contact terms between the hidden sector field X and the 5D MSSM multiplets in the bulk. We introduce a contact term between the gauge multiplets and X of the form

$$S_5^{Xg} = \int d^4x \int dy \left\{ \int d^2\theta d_a \frac{X}{M_5^2} \text{Tr}[\tilde{W}^\alpha \tilde{W}_\alpha] + \text{H.c.} \right\} \delta(y), \quad (18)$$

where the original \tilde{W}^α has a mass dimension of 2 before normalizing V . The contact terms between the Higgs hypermultiplets and the hidden sector field X are written as

$$S_5^{Xh} = \int d^4x \int dy \times \int d^4\theta \left\{ \left[d_\mu \frac{X^\dagger}{M_5^2} H_u^h H_d^h + d_{B_\mu} \frac{X^\dagger X}{M_5^3} H_u^h H_d^h + \text{H.c.} \right] + \left[d_A^{H_u} \frac{X + X^\dagger}{M_5^2} H_u^{h\dagger} H_u^h + d_m^{H_u} \frac{X^\dagger X}{M_5^3} H_u^{h\dagger} H_u^h + d_A^{H_d} \frac{X + X^\dagger}{M_5^2} H_d^{h\dagger} H_d^h + d_m^{H_d} \frac{X^\dagger X}{M_5^3} H_d^{h\dagger} H_d^h \right] \right\} \delta(y), \quad (19)$$

and those between the matter hypermultiplets and X are

$$S_5^{Xm} = \int d^4x \int dy \int d^4\theta \times \left[(d_A^\Phi)_{ij} \frac{X + X^\dagger}{M_5^2} \Phi_i^\dagger \Phi_j + (d_m^\Phi)_{ij} \frac{X^\dagger X}{M_5^3} \Phi_i^\dagger \Phi_j \right] \delta(y), \quad (20)$$

where Φ stands for $\{Q^h, U^h, D^h, L^h, E^h\}$ hypermultiplets, and $\{i, j\}$ are generation indices. We can also introduce contact terms in the superpotential as follows:

$$S_5^{Xa} = \int d^4x \int dy \int d^2\theta \left\{ \frac{(a_u)_{ij}}{M_5^{5/2}} X H_u^h Q_i^h U_j^h + \frac{(a_d)_{ij}}{M_5^{5/2}} X H_d^h Q_i^h D_j^h + \frac{(a_e)_{ij}}{M_5^{5/2}} X H_d^h L_i^h E_j^h + \text{H.c.} \right\} \delta(y). \quad (21)$$

The SUSY breaking by $\langle F_X \rangle$ induces the soft SUSY breaking terms in the MSSM through the above contact terms at the effective 4D cutoff scale $M_{\text{cut}} = M_4 \omega$. The gaugino masses are given by

$$M_a \simeq -\sqrt{3} \left(\frac{d_a}{2\pi R M_4} \right) g_a^2 m_{3/2}, \quad (a = 1, 2, 3), \quad (22)$$

where g_a is the SM gauge coupling, and the gravitino mass $m_{3/2}$ is given by

$$m_{3/2} = \frac{\langle F_X \rangle}{\sqrt{3} M_4}. \quad (23)$$

Note that the gaugino mass is suppressed by the so-called volume suppression factor of $1/(2\pi R M_4)$. Soft masses of the Higgs sector can be obtained from Eq. (19),

$$m_{H_u}^2 = 3[-d_m^{H_u} + (d_A^{H_u})^2 \mathcal{C}_{H_u}^2] \mathcal{C}_{H_u}^2 m_{3/2}^2, \quad (24)$$

$$m_{H_d}^2 = 3[-d_m^{H_d} + (d_A^{H_d})^2 \mathcal{C}_{H_d}^2] \mathcal{C}_{H_d}^2 m_{3/2}^2, \quad (25)$$

$$B_\mu = 3d_{B_\mu} \mathcal{C}_{H_u} \mathcal{C}_{H_d} m_{3/2}^2 - \sqrt{3}(d_A^{H_u} \mathcal{C}_{H_u}^2 + d_A^{H_d} \mathcal{C}_{H_d}^2) \mu m_{3/2}. \quad (26)$$

In our R -charge assignment, the μ -term is forbidden, but it is generated through the SUSY breaking [23],

$$\mu = \sqrt{3} d_\mu \mathcal{C}_{H_u} \mathcal{C}_{H_d} m_{3/2}. \quad (27)$$

Scalar soft masses of sparticles are generated from Eq. (20),

$$(m_\Phi)_{ij}^2 = 3 \left[-(d_m^\Phi)_{ij} + \sum_{n=1}^3 (d_A^\Phi)_{in} (d_A^\Phi)_{nj} \mathcal{C}_{\Phi_n}^2 \right] \mathcal{C}_{\Phi_i} \mathcal{C}_{\Phi_j} m_{3/2}^2, \quad (28)$$

where Φ stands for Q, U, D, L, E , and $i, j = \{1, 2, 3\}$. Last but not least, the trilinear coupling A -terms arise from Eq. (21),

$$(A_u)_{ij} = \frac{\sqrt{3} m_{3/2}}{(Y_u)_{ij}} \left[(a_u)_{ij} \mathcal{C}_{H_u} \mathcal{C}_{Q_i} \mathcal{C}_{Q_j} - d_A^{H_u} (Y_u)_{ij} \mathcal{C}_{H_u}^2 - \sum_{n=1}^3 (d_A^Q)_{ni} (Y_u)_{nj} \mathcal{C}_{Q_n} \mathcal{C}_{Q_i} - \sum_{n=1}^3 (d_A^U)_{nj} (Y_u)_{in} \mathcal{C}_{U_n} \mathcal{C}_{U_j} \right], \quad (29)$$

$$(A_d)_{ij} = \frac{\sqrt{3} m_{3/2}}{(Y_d)_{ij}} \left[(a_d)_{ij} \mathcal{C}_{H_d} \mathcal{C}_{Q_i} \mathcal{C}_{D_j} - d_A^{H_d} (Y_d)_{ij} \mathcal{C}_{H_d}^2 - \sum_{n=1}^3 (d_A^Q)_{ni} (Y_d)_{nj} \mathcal{C}_{Q_n} \mathcal{C}_{Q_i} - \sum_{n=1}^3 (d_A^D)_{nj} (Y_d)_{in} \mathcal{C}_{D_n} \mathcal{C}_{D_j} \right], \quad (30)$$

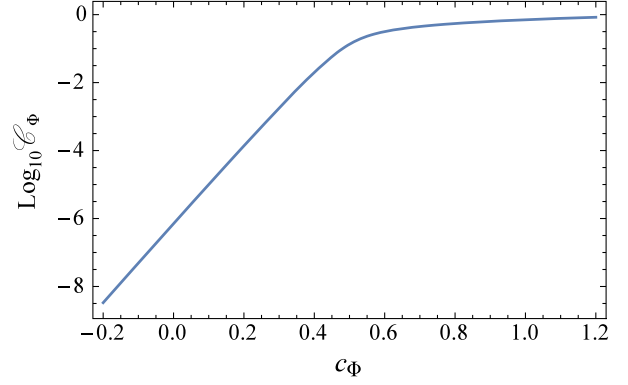


FIG. 1. Behavior of the geometrical coefficient as a function of the bulk mass parameter: $\mathcal{C}_\Phi = \sqrt{\frac{(1-2c_\Phi)}{2(\omega^{-1+2c_\Phi}-1)}}$. The plot is a demonstration for the case with $\omega = 10^{-12}$.

$$(A_e)_{ij} = \frac{\sqrt{3} m_{3/2}}{(Y_e)_{ij}} \left[(a_e)_{ij} \mathcal{C}_{H_d} \mathcal{C}_{L_i} \mathcal{C}_{E_j} - d_A^{H_d} (Y_e)_{ij} \mathcal{C}_{H_d}^2 - \sum_{n=1}^3 (d_A^L)_{ni} (Y_e)_{nj} \mathcal{C}_{L_n} \mathcal{C}_{L_i} - \sum_{n=1}^3 (d_A^E)_{nj} (Y_e)_{in} \mathcal{C}_{E_n} \mathcal{C}_{E_j} \right]. \quad (31)$$

To avoid the SUSY flavor changing neutral currents (FCNCs), we assume that the couplings d_m^Φ , d_A^Φ , a_u , a_d , a_e are all flavor diagonal and, in particular, flavor universal for the first two generations.

All terms induced by the SUSY breaking are controlled by the gravitino mass and the coupling constants for the contact interactions between the MSSM multiplets and the hidden sector field X . In addition, the warped background geometry plays a crucial role in determining the size of the parameters. The gaugino masses are volume suppressed from the gravitino mass. The scalar masses squared, the A -terms and the μ -term are controlled by the geometrical coefficients of bulk hypermultiplets \mathcal{C}_Φ . In Fig. 1, we show a geometrical coefficient \mathcal{C}_Φ as a function of the bulk mass parameter c_Φ . In this example, we have set $\omega = 10^{-12}$, which generates the low cutoff scale of $\mathcal{O}(10^6)$ GeV.⁴ For $c_\Phi = 1/2$, the wave function is independent of y , and $\mathcal{C}_\Phi = \mathcal{O}(0.1)$ in this case corresponds to the volume suppression factor. As $c_\Phi > 1/2$ increases, the wave function tends to localize towards the UV brane, and hence \mathcal{C}_Φ is approaching 1. On the other hand, as $c_\Phi < 1/2$ decreases, the wave function tends to localize towards the IR brane, and \mathcal{C}_Φ is exponentially suppressed. Therefore, with a suitable choice of the bulk parameters, we can easily achieve a hierarchy

⁴Basically, the choice of ω is arbitrary. But we have found that high cutoff scales result in tachyonic staus. Therefore, the choice of the cutoff scale of $\mathcal{O}(10^6)$ GeV is preferable.

between soft SUSY breaking parameters. From Eqs. (22)–(31), we see that the maximum values of the soft parameters are of $\mathcal{O}(m_{3/2})$ when the coupling constants are of the order 1. For $\omega \ll 1$, the gaugino masses are roughly an order of magnitude smaller than the gravitino mass. Thus, the model predicts the dark matter candidate to be the lightest neutralino. We find that soft parameters for the scalars localized around the UV brane are of $\mathcal{O}(m_{3/2})$, while they can be much smaller for the scalars localized around the IR brane. Interestingly, once all the couplings between the 5D MSSM multiplet and the hidden sector field are set to be universal, for instance of $\mathcal{O}(1)$, the diversity of soft masses and couplings of the 4D MSSM can

be derived from the warped geometry with appropriate localization.

III. BENCHMARK PARTICLE MASS SPECTRUM

In this section, we investigate realistic particle mass spectra that satisfy all phenomenological constraints. As we discussed in the previous section, a suitable choice of the bulk mass parameters can naturally generate a hierarchy between sparticle masses. In the following analysis, we consider the 4D effective MSSM with the inputs of soft SUSY breaking parameters at the 4D effective cutoff scale of $M_{\text{cut}} \simeq M_4 \omega$. The low-energy mass spectrum is obtained

TABLE II. Benchmark particle mass spectra in GeV units for $M_{\text{cut}} = 10^5$ GeV. Input soft masses for the first two generation squarks, and the third generation slepton and squark, are fixed as $m_q^0 = 8500$ GeV, $m_{l3}^0 = 800$ GeV, and $m_{q3}^0 = 9500$ GeV. Other parameters including the gaugino sector input m_g , the soft masses for the first two generation sleptons m_l^0 and for two Higgs doublets m_h^0 , the universal trilinear coupling A_0 , and $\tan \beta$ are allowed to vary in this table.

M_{cut}	10^5	10^5	10^5	10^5
m_g	1120	1178	1330	1615
m_l^0, m_q^0	306.0, 8500	312.1, 8500	331.0, 8500	370.1, 8500
m_{l3}^0, m_{q3}^0	800, 9500	800, 9500	800, 9500	800, 9500
m_h^0	1000	2500	2600	2800
A_0	-6800	-2200	-1000	0
$\tan \beta$	10	20	30	40
h^0	125.27	125.28	125.29	125.22
H^0, A^0	3716	3307	3057	2766
H^\pm	3717	3308	3059	2767
\tilde{g}	1405	1469	1639	1953
$\tilde{\chi}_{1,2}^0$	245, 484	257, 507	290, 571	352, 689
$\tilde{\chi}_{3,4}^0$	3600, 3601	2426, 2427	2278, 2279	2029, 2030
$\tilde{\chi}_{1,2}^\pm$	484, 3602	507, 2428	571, 2280	689, 2031
$\tilde{u}, \tilde{c}_{L,R}$	8572, 8560	8573, 8561	8577, 8565	8585, 8573
$\tilde{d}, \tilde{s}_{L,R}$	8572, 8559	8573, 8560	8577, 8564	8585, 8572
$\tilde{t}_{1,2}$	9018, 9317	9112, 9339	9124, 9317	9130, 9283
$\tilde{b}_{1,2}$	9307, 9555	9336, 9525	9316, 9474	9281, 9401
$\tilde{\nu}_L^{\tau,\mu}$	253, 253	265, 265	298, 298	358, 358
$\tilde{e}, \tilde{\mu}_L$	265, 265	277, 277	308, 308	367, 367
$\tilde{e}, \tilde{\mu}_R$	335, 335	337, 337	354, 354	393, 392
$\tilde{\nu}_L^e$	765	766	760	751
$\tilde{\tau}_{1,2}$	726, 819	716, 827	677, 835	633, 833
Δa_μ	2.58×10^{-9}	2.94×10^{-9}	3.23×10^{-9}	2.62×10^{-9}
$\text{BR}(b \rightarrow s + \gamma)$	3.34×10^{-4}	3.34×10^{-4}	3.35×10^{-4}	3.36×10^{-4}
$\text{BR}(B_s \rightarrow \mu^+ + \mu^-)$	3.08×10^{-9}	3.07×10^{-9}	3.04×10^{-9}	2.98×10^{-9}
$\frac{\text{BR}^{\text{exp}}(B_u \rightarrow \tau + \nu_\tau)}{\text{BR}^{\text{SM}}(B_u \rightarrow \tau + \nu_\tau)}$	1.00	9.98×10^{-1}	9.95×10^{-1}	9.89×10^{-1}
$\text{BR}(D_s \rightarrow \tau + \nu_\tau)$	5.17×10^{-2}	5.17×10^{-2}	5.17×10^{-2}	5.17×10^{-2}
$\text{BR}(D_s \rightarrow \mu + \nu_\mu)$	5.33×10^{-3}	5.33×10^{-3}	5.33×10^{-3}	5.33×10^{-3}
R_{l23}	1.000	1.000	1.000	1.000
Ωh^2	0.119	0.119	0.119	0.119
$\sigma_{\text{SI}}^{\chi^0-p}$ (pb)	9.35×10^{-13}	1.31×10^{-12}	1.46×10^{-12}	2.75×10^{-12}
$\sigma_{\text{SD}}^{\chi^0-p}$ (pb)	6.22×10^{-10}	3.88×10^{-9}	5.17×10^{-9}	8.74×10^{-9}

through the renormalization group (RG) evolutions. We employ the SOFTSUSY package (version 3.6.2) [24] to numerically solve the RG equations. With the output at low energies, other physical observables and constraints are computed by using the MicrOMEGAs package (version 4.2.3) [25].

Regarding the inputs for the MSSM gaugino masses, we simply assume the universal couplings, $d_1 = d_2 = d_3$. Since the ratio $m_g = M_a/g_a^2$ is RG invariant at the one-loop level, the resultant mass ratio among the gauginos is the same as those in the constrained MSSM when m_g is set as a common input for the gaugino sector at the cutoff scale. Hence the bino and wino are lighter than the gluino. To avoid the severe experimental constraints on the SUSY FCNCs for the first two generations, we assume that the couplings of the hidden sector with the first two generation matter fields are flavor blind. In order to simplify our

analysis, we set the universal soft mass inputs for the two Higgs doublets (m_h^0), and the sleptons and squarks in the first two generations (m_l^0, m_q^0). The other free parameters in our analysis are the universal A-term A_0 at M_{cut} and $\tan\beta$. We choose $\text{sign}(\mu) = +1$ to yield a positive Δa_μ that can fill up the discrepancy between the experimental value and the SM prediction.

In our study, we consider various phenomenological constraints. We employ the combined result for the Higgs boson mass measured by the ATLAS and the CMS collaborations [4]. The lower mass bounds on squarks and gluino in the simplified model [26] are taken into account as a reference. As a motivation of this paper, the benchmark points are chosen such that the muon anomalous magnetic moment $a_\mu = \frac{g_\mu - 2}{2}$ satisfies the current experimental value [1,2]. Other constraints are from the branching ratios of rare decay processes: $b \rightarrow s + \gamma$ [27],

TABLE III. Same as Table II, but for $M_{\text{cut}} = 10^6$ GeV.

M_{cut}	10^6	10^6	10^6	10^6
m_g	1114	1300	1350	1400
m_l^0, m_q^0	349.4, 8500	365.3, 8500	370.5, 8500	374.0, 8500
m_{l3}^0, m_{q3}^0	800, 9500	800, 9500	800, 9500	800, 9500
m_h^0	3200	3500	4000	4300
A_0	-6450	-2600	-2000	-400
$\tan\beta$	10	20	30	40
h^0	125.00	125.16	125.21	125.22
H^0, A^0	4834	4343	4036	3541
H^\pm	4835	4344	4036	3542
\tilde{g}	1400	1606	1661	1715
$\tilde{\chi}_{1,2}^0$	243, 480	283, 557	294, 578	304, 596
$\tilde{\chi}_{3,4}^0$	3675, 3676	2945, 2946	2239, 2240	1558, 1560
$\tilde{\chi}_{1,2}^\pm$	480, 3676	558, 2947	578, 2241	596, 1561
$\tilde{u}, \tilde{c}_{L,R}$	8557, 8546	8564, 8554	8566, 8556	8569(8), 8558
$\tilde{d}, \tilde{s}_{L,R}$	8557, 8546	8564, 8553	8567, 8555	8569, 8557
$\tilde{t}_{1,2}$	8666, 9130	8788, 9163	8784, 9113	8791, 9045
$\tilde{b}_{1,2}$	9125, 9535	9161, 9489	9111, 9396	9044, 9259
$\tilde{\nu}_L^{\tau,\mu}$	252, 252	291, 291	302, 301	312, 311
$\tilde{e}, \tilde{\mu}_L$	264, 264	302, 302	312, 312	322, 321
$\tilde{e}, \tilde{\mu}_R$	368, 368	382, 381	386, 385	387, 386
$\tilde{\nu}_L^e$	736	729	680	631
$\tilde{\tau}_{1,2}$	699, 797	657, 799	554, 745	440, 681
Δa_μ	2.31×10^{-9}	2.58×10^{-9}	2.88×10^{-9}	2.84×10^{-9}
$\text{BR}(b \rightarrow s + \gamma)$	3.33×10^{-4}	3.33×10^{-4}	3.33×10^{-4}	3.34×10^{-4}
$\text{BR}(B_s \rightarrow \mu^+ + \mu^-)$	3.08×10^{-9}	3.07×10^{-9}	3.06×10^{-9}	3.02×10^{-9}
$\frac{\text{BR}^{\text{exp}}(B_s \rightarrow \tau + \nu_\tau)}{\text{BR}^{\text{SM}}(B_s \rightarrow \tau + \nu_\tau)}$	1.00	9.99×10^{-1}	9.97×10^{-1}	9.93×10^{-1}
$\text{BR}(D_s \rightarrow \tau + \nu_\tau)$	5.17×10^{-2}	5.17×10^{-2}	5.17×10^{-2}	5.17×10^{-2}
$\text{BR}(D_s \rightarrow \mu + \nu_\mu)$	5.33×10^{-3}	5.33×10^{-3}	5.33×10^{-3}	5.33×10^{-3}
R_{l23}	1.000	1.000	1.000	1.000
Ωh^2	0.119	0.119	0.119	0.119
σ_{SI}^{-p} (pb)	8.82×10^{-13}	7.78×10^{-13}	1.44×10^{-12}	4.96×10^{-12}
σ_{SD}^{-p} (pb)	5.65×10^{-10}	1.66×10^{-9}	5.59×10^{-9}	2.66×10^{-8}

$B_s \rightarrow \mu^+ + \mu^-$ [28], $B \rightarrow \tau + \nu_\tau$ [29], $D_s \rightarrow \tau + \nu_\tau$ [30],
 $D_s \rightarrow \mu + \nu_\mu$ [30], and the kaon decay parameter [31],

$$R_{l23} = \left| \frac{V_{us}(K_{l2})}{V_{us}(K_{l3})} \times \frac{V_{ud}(0^+ \rightarrow 0^+)}{V_{ud}(\pi_{l2})} \right|, \quad (32)$$

where the CKM matrix elements, V_{us} and V_{ud} , are measured from the corresponding three-body semileptonic kaon decay (K_{l3}), two-body leptonic kaon and pion decay (K_{l2}, π_{l2}), and superallowed nuclear beta decay ($0^+ \rightarrow 0^+$). The constraints that we employ are listed below.

$$m_h = 125.09 \pm 0.21(\text{stat.}) \pm 0.11(\text{syst.}) \text{ GeV}, \quad (33)$$

$$m_{\tilde{g}} \gtrsim 1.4 \text{ TeV}, \quad (34)$$

$$\Delta a_\mu = a_\mu^{\text{exp}} - a_\mu^{\text{SM}} = (28.6 \pm 8.0) \times 10^{-10}, \quad (35)$$

$$2.99 \times 10^{-4} < \text{BR}(b \rightarrow s + \gamma) < 3.87 \times 10^{-4}, \quad (2\sigma) \quad (36)$$

$$2.1 \times 10^{-9} < \text{BR}(B_s \rightarrow \mu^+ + \mu^-) < 4.0 \times 10^{-9}, \quad (1\sigma)$$

$$(37)$$

$$0.15 < \frac{\text{BR}^{\text{exp}}(B_u \rightarrow \tau + \nu_\tau)}{\text{BR}^{\text{SM}}(B_u \rightarrow \tau + \nu_\tau)} < 2.41, \quad (3\sigma) \quad (38)$$

TABLE IV. Benchmark particle mass spectra in GeV units for $M_{\text{cut}} = 10^7$ GeV. Input soft masses for the first two generation squarks, and the third generation slepton and squark, are fixed as $m_q^0 = 8500$ GeV, $m_{l3}^0 = 800$ GeV, and $m_{q3}^0 = 9500$ GeV. In the last column, the input for the third generation slepton mass is taken to be a bit larger, $m_{l3}^0 = 1000$ GeV.

M_{cut}	10^7	10^7	10^7	10^7
$m_{\tilde{g}}$	1280	1260	1230	1320
m_l^0, m_q^0	388.9, 8500	397.6, 8500	396.4, 8500	400.6, 8500
m_{l3}^0, m_{q3}^0	800, 9500	800, 9500	800, 9500	1000, 9500
m_h^0	1000	4500	5000	5100
A_0	-7500	-4000	-500	0
$\tan \beta$	10	20	30	40
h^0	125.02	125.13	125.08	125.11
H^0, A^0	5611	5186	4723	4142
H^\pm	5612	5187	4724	4143
\tilde{g}	1589	1563	1524	1624
$\tilde{\chi}_{1,2}^0$	280, 551	275, 540	267, 524	286, 561
$\tilde{\chi}_{3,4}^0$	5548, 5548	3113, 3114	1966, 1968	1698, 1700
$\tilde{\chi}_{1,2}^\pm$	551, 5549	540, 3114	524, 1969	561, 1701
$\tilde{u}, \tilde{c}_{L,R}$	8553, 8544	8551, 8542	8549, 8541	8555(4), 8546
$\tilde{d}, \tilde{s}_{L,R}$	8553, 8542	8551, 8541	8550, 8540	8555, 8545(4)
$\tilde{t}_{1,2}$	8377, 8984	8435, 8972	8476, 8926	8478, 8830
$\tilde{b}_{1,2}$	8979, 9525	8970, 9448	8924, 9323	8829, 9136
$\tilde{\nu}_L^{e,\mu}$	288, 288	283, 282	276, 275	295, 293
$\tilde{e}, \tilde{\mu}_L$	298, 298	294, 293	287, 286	306, 304
$\tilde{e}, \tilde{\mu}_R$	393, 392	404, 403	400, 399	403, 401
$\tilde{\nu}_L^\tau$	713	651	598	745
$\tilde{\tau}_{1,2}$	643, 788	536, 717	432, 648	471, 778
Δa_μ	2.48×10^{-9}	2.69×10^{-9}	2.84×10^{-9}	3.13×10^{-9}
$\text{BR}(b \rightarrow s + \gamma)$	3.32×10^{-4}	3.32×10^{-4}	3.33×10^{-4}	3.33×10^{-4}
$\text{BR}(B_s \rightarrow \mu^+ + \mu^-)$	3.08×10^{-9}	3.08×10^{-9}	3.06×10^{-9}	3.03×10^{-9}
$\frac{\text{BR}^{\text{exp}}(B_u \rightarrow \tau + \nu_\tau)}{\text{BR}^{\text{SM}}(B_u \rightarrow \tau + \nu_\tau)}$	1.00	9.99×10^{-1}	9.98×10^{-1}	9.95×10^{-1}
$\text{BR}(D_s \rightarrow \tau + \nu_\tau)$	5.17×10^{-2}	5.17×10^{-2}	5.17×10^{-2}	5.17×10^{-2}
$\text{BR}(D_s \rightarrow \mu + \nu_\mu)$	5.33×10^{-3}	5.33×10^{-3}	5.33×10^{-3}	5.33×10^{-3}
R_{l23}	1.000	1.000	1.000	1.000
Ωh^2	0.119	0.119	0.119	0.119
$\sigma_{\text{SI}}^{\chi-P}(\text{pb})$	3.49×10^{-13}	6.34×10^{-13}	1.89×10^{-12}	3.17×10^{-12}
$\sigma_{\text{SD}}^{\chi-P}(\text{pb})$	5.81×10^{-11}	1.29×10^{-9}	9.70×10^{-9}	1.83×10^{-8}

$$5.07 \times 10^{-2} < \text{BR}(D_s \rightarrow \tau + \nu_\tau) < 6.03 \times 10^{-2}, \quad (2\sigma) \quad (39)$$

$$5.31 \times 10^{-3} < \text{BR}(D_s \rightarrow \mu + \nu_\mu) < 5.81 \times 10^{-3}, \quad (1\sigma) \quad (40)$$

$$R_{l23} = 1.004 \pm 0.007. \quad (41)$$

Since the LHC constraints require that sparticles must be heavy, their contributions to the precision electroweak observables are negligibly small [32].

Assuming R-parity conservation, the lightest neutralino is a primary candidate of the cold dark matter. Beside the above constraints, we also consider the cosmological constraint on the neutralino dark matter relic abundance. Here we apply the result by the Planck satellite experiment [33],

$$\Omega h^2 = 0.1188 \pm 0.0010 \quad (68\% \text{ C.L.}). \quad (42)$$

Finally, the constraints from the results of the direct and indirect dark matter searches are taken into account. The most stringent upper limit on the spin-independent cross section of the neutralino dark matter with nucleon has been reported by the LUX experiment [34], while the IceCube experiment has set the most severe upper limit on the spin-dependent cross section between the neutralino dark matter and nucleon [35],

$$\sigma_{\text{SI}}^{\chi-P} \lesssim 7 \times 10^{-9} \text{ pb}(90\% \text{ C.L.}), \quad \text{for } m_{\text{WIMP}} \approx 600 \text{ GeV}, \quad (43)$$

$$\sigma_{\text{SD}}^{\chi-P} \lesssim 10^{-4} \text{ pb}(90\% \text{ C.L.}), \quad \text{for } m_{\text{WIMP}} \approx 150\text{--}600 \text{ GeV}. \quad (44)$$

The benchmark mass spectra along with the observables satisfying the above phenomenological constraints are shown in Tables II–IV for three cutoff scales, $M_{\text{cut}} = 10^5, 10^6, \text{ and } 10^7 \text{ GeV}$, respectively. In the tables, $m_q^0 = 8500 \text{ GeV}$ is the common input soft mass for all squarks in the first two generations, while the common masses for the third generation sleptons and squarks are fixed to be $m_{l3}^0 = 800 \text{ GeV}$ and $m_{q3}^0 = 9500 \text{ GeV}$, respectively. The choice of m_{l3}^0 for the benchmark point of the last column of Table IV is a bit larger. The other four input parameters at the cutoff scale $\{m_g, m_l^0, m_h^0, A_0\}$ are chosen so as to satisfy the four most important constraints: the Higgs boson mass, gluino mass, the muon anomalous magnetic dipole moment, and the dark matter relic density. The mass hierarchy between squarks and sleptons/gauginos is crucial to reproduce $m_h \approx 125 \text{ GeV}$ and $\Delta a_\mu = \mathcal{O}(10^{-9})$

simultaneously. The benchmark points satisfy all these phenomenological constraints.

The dark matter neutralinos in the benchmarks are all binolike. Since we have chosen m_{l3}^0 larger than m_l^0 to avoid the stau being the lightest sparticle (LSP), the next-to-LSP (NLSP) is the muon sneutrino that is almost degenerate with the electron sneutrino. The right dark matter relic abundance is achieved through the coannihilation processes between the neutralino LSP and the electron/muon sneutrinos, which is ensured by a correlation between the free inputs m_g and m_l^0 . As can be seen in Eq. (2), the sparticle contribution to the muon $g-2$ is proportional to $\tan\beta$. Hence, as the input of $\tan\beta$ is raised, the inputs of m_g and m_l^0 are increased to satisfy the constraint from the muon $g-2$.

The cutoff scales in all the tables are just about a few orders of magnitude higher than typical squark masses. Therefore, in our model, the distance of the RG evolutions of soft SUSY breaking parameters are very short compared to, for example, the constrained MSSM, and hence the RG evolution effects are much less. In addition, the inputs of the scalar squared masses are nonuniversal. In the slepton sector, we can see from the tables that the slepton masses at low energy are smaller than the corresponding inputs at the boundary. This is the effect from the higher order corrections in the RG equations with the hierarchically large inputs of the squark masses. Not as in the constrained MSSM, the left-handed sleptons of the first two generations become lighter than right-handed ones. As a consequence, the NLSP in the provided spectra is muon sneutrino.

IV. CONCLUSIONS

In order to reconcile the Higgs boson mass $m_h \approx 125 \text{ GeV}$ and the discrepancy of the muon anomalous magnetic dipole moment $\Delta a_\mu \sim 10^{-9}$, a hierarchical mass splitting between squarks and sleptons/gauginos is usually necessary. In this paper, we have presented a 5D MSSM in the RS warped background metric with the warp factor $\omega \ll 1$. All the MSSM multiplets reside in the bulk, while the SUSY is broken on the UV brane where a hidden chiral field is localized. The zero modes of the 5D MSSM fields are identified as the MSSM fields in the 4D effective theory. The SUSY breaking mediation to the MSSM sector is controlled by how much the MSSM sparticles in the bulk overlap with the hidden field on the UV brane. Since the gravitino is localizing around the UV brane, the SUSY breaking parameters in the MSSM are characterized by the gravitino mass $m_{3/2}$ and geometrical factors corresponding to the zero-mode configurations. The gaugino masses are volume suppressed $\sim 0.1m_{3/2}$. Squarks are localized around the UV brane with a bulk mass parameter $> 1/2$, while leptons acquire much smaller

masses with a suitable choice of the bulk mass parameter $< 1/2$. Interestingly, assuming a common coupling between the hidden sector and the 5D MSSM sector, the diversity of the 4D MSSM soft terms can be derived from the universality of the underlying theory with the warped geometry and appropriate localization. In our setup, a factor deference between bulk mass parameters, which are the original parameters in the model, results in a hierarchy because of the warped metric. With the hierarchical mass spectrum generated by the warped geometry, we have demonstrated with the benchmarks that not only $m_h \approx 125$ GeV and $\Delta a_\mu \sim 10^{-9}$ can be reconciled, but also various phenomenological constraints such as the right abundance of the neturinalino dark matter, the SUSY FCNC constraints, and the LHC

bounds on sparticle masses are satisfied. In the benchmark points, squarks are too heavy to be produced at the LHC, while sleptons, light charginos, and neutralinos can be explored at the LHC Run-2 in the future.

ACKNOWLEDGMENTS

H. M. T. thanks the Department of Physics and Astronomy at the University of Alabama for hospitality during his visit. The work of N. O. is supported in part by the United States Department of Energy (Award No. DE-SC0013680). The work of H. M. T. is supported in part by the Vietnam National Foundation for Science and Technology Development (NAFOSTED) under Grant No. 103.01-2014.22.

-
- [1] G. W. Bennett *et al.* (Muon g-2 Collaboration), *Phys. Rev. D* **73**, 072003 (2006).
- [2] K. Hagiwara, R. Liao, A. D. Martin, D. Nomura, and T. Teubner, *J. Phys. G* **38**, 085003 (2011); M. Davier, A. Hoecker, B. Malaescu, and Z. Zhang, *Eur. Phys. J. C* **71**, 1515 (2011); **72**, 1874 (2012); M. Benayoun, P. David, L. DelBuono, and F. Jegerlehner, *Eur. Phys. J. C* **73**, 2453 (2013); T. Blum, A. Denig, I. Logashenko, E. de Rafael, B. Lee Roberts, T. Teubner, and G. Venanzoni, arXiv: 1311.2198.
- [3] See, for example, T. Moroi, *Phys. Rev. D* **53**, 6565 (1996); **56**, 4424 (1997); S. Komine, T. Moroi, and M. Yamaguchi, *Phys. Lett. B* **506**, 93 (2001); N. Okada and H. M. Tran, *Phys. Rev. D* **83**, 053001 (2011); H. M. Tran, T. Kon, and Y. Kurihara, *Mod. Phys. Lett. A* **26**, 949 (2011); J. L. Evans, M. Ibe, S. Shirai, and T. T. Yanagida, *Phys. Rev. D* **85**, 095004 (2012); J. J. Cao, Z. X. Heng, J. M. Yang, Y. M. Zhang, and J. Y. Zhu, *J. High Energy Phys.* 03 (2012) 086; K. Harigaya, T. T. Yanagida, and N. Yokozaki, *Phys. Rev. D* **91**, 075010 (2015).
- [4] G. Aad *et al.* (ATLAS and CMS Collaborations), *Phys. Rev. Lett.* **114**, 191803 (2015).
- [5] H. E. Haber, R. Hempfling, and A. H. Hoang, *Z. Phys. C* **75**, 539 (1997); A. Djouadi, *Phys. Rep.* **459**, 1 (2008).
- [6] J. Cao, Z. Heng, D. Li, and J. M. Yang, *Phys. Lett. B* **710**, 665 (2012); N. Okada and H. M. Tran, *Phys. Rev. D* **87**, 035024 (2013).
- [7] T. Moroi, *Phys. Rev. D* **53**, 6565 (1996); **56**, 4424 (1997); S. P. Martin and J. D. Wells, *Phys. Rev. D* **64**, 035003 (2001); G. F. Giudice, P. Paradisi, A. Strumia, and A. Strumia, *J. High Energy Phys.* 10 (2012) 186.
- [8] R. Sato, K. Tobioka, and N. Yokozaki, *Phys. Lett. B* **716**, 441 (2012); M. Ibe, S. Matsumoto, T. T. Yanagida, and N. Yokozaki, *J. High Energy Phys.* 03 (2013) 078; M. Endo, K. Hamaguchi, S. Iwamoto, and T. Yoshinaga, *J. High Energy Phys.* 01 (2014) 123; M. Ibe, T. T. Yanagida, and N. Yokozaki, *J. High Energy Phys.* 08 (2013) 067; S. Akula and P. Nath, *Phys. Rev. D* **87**, 115022 (2013); M. Endo, K. Hamaguchi, T. Kitahara, and T. Yoshinaga, *J. High Energy Phys.* 11 (2013) 013; M. A. Ajaib, I. Gogoladze, Q. Shafi, and C. S. Un, *J. High Energy Phys.* 05 (2014) 079; I. Gogoladze, F. Nasir, Q. Shafi, and C. S. Un, *Phys. Rev. D* **90**, 035008 (2014); S. P. Das, M. Guchait, and D. P. Roy, *Phys. Rev. D* **90**, 055011 (2014); M. Badziak, Z. Lalak, M. Lewicki, M. Olechowski, and S. Pokorski, *J. High Energy Phys.* 03 (2015) 003; M. Adeel Ajaib, I. Gogoladze, and Q. Shafi, *Phys. Rev. D* **91**, 095005 (2015); K. Kowalska, L. Roszkowski, E. M. Sessolo, and A. J. Williams, *J. High Energy Phys.* 06 (2015) 020; F. Wang, W. Wang, and J. M. Yang, *J. High Energy Phys.* 06 (2015) 079; F. Wang, W. Wang, J. M. Yang, and Y. Zhang, *J. High Energy Phys.* 07 (2015) 138; D. Chowdhury and N. Yokozaki, *J. High Energy Phys.* 08 (2015) 111; M. A. Ajaib, B. Dutta, T. Ghosh, I. Gogoladze, and Q. Shafi, *Phys. Rev. D* **92**, 075033 (2015); I. Gogoladze, Q. Shafi, and C. S. Un, *Phys. Rev. D* **92**, 115014 (2015); F. Wang, L. Wu, J. M. Yang, and M. Zhang, *Phys. Lett. B* **759**, 191 (2016); T. Li, S. Raza, and K. Wang, *Phys. Rev. D* **93**, 055040 (2016).
- [9] N. Okada, S. Raza, and Q. Shafi, *Phys. Rev. D* **90**, 015020 (2014); K. S. Babu, I. Gogoladze, Q. Shafi, C. S. Un, T. Li, S. Raza, and K. Wang, *Phys. Rev. D* **93**, 055040 (2016).
- [10] K. S. Babu, I. Gogoladze, Q. Shafi, and C. S. Un, *Phys. Rev. D* **90**, 116002 (2014).
- [11] A. Abdalgabar, A. S. Cornell, A. Deandrea, and M. McGarrie, *J. High Energy Phys.* 07 (2014) 158.
- [12] M. Endo, K. Hamaguchi, S. Iwamoto, and N. Yokozaki, *Phys. Rev. D* **84**, 075017 (2011); **85**, 095012 (2012); S. Iwamoto, *AIP Conf. Proc.* **1467**, 57 (2012); S. Mohanty, S. Rao, and D. P. Roy, *J. High Energy Phys.* 09 (2013) 027; J. H. Huh and B. Kyae, *Phys. Lett. B* **726**, 729 (2013); M. Arana-Catania, S. Heinemeyer, and M. J. Herrero, *Phys. Rev. D* **90**, 075003 (2014); S. Iwamoto, T. T. Yanagida, and N. Yokozaki, *Prog. Theor. Exp. Phys.* **2015**, 073B01 (2015); L. Calibbi, P. Paradisi, and R. Ziegler, *Eur. Phys. J. C* **74**,

- 3211 (2014); D. T. Binh, D. T. Huong, and H. N. Long, *Zh. Eksp. Teor. Fiz.* **148**, 1115 (2015) [*J. Exp. Theor. Phys.* **121**, 976 (2015)]; K. Harigaya, T. T. Yanagida, and N. Yokozaki, *Phys. Rev. D* **92**, 035011 (2015); G. Bélanger, J. Da Silva, U. Laa, and A. Pukhov, *J. High Energy Phys.* **09** (2015) 151; S. Khalil and C. S. Un, arXiv:1509.05391; A. Chakraborty and S. Chakraborty, *Phys. Rev. D* **93**, 075035 (2016); F. V. Flores-Baez, M. G. Bock, and M. Mondragon, arXiv:1512.00902; M. Nishida and K. Yoshioka, arXiv:1605.06675.
- [13] L. Randall and R. Sundrum, *Phys. Rev. Lett.* **83**, 3370 (1999).
- [14] W. D. Goldberger and M. B. Wise, *Phys. Rev. D* **60**, 107505 (1999); S. Chang, J. Hisano, H. Nakano, N. Okada, and M. Yamaguchi, *Phys. Rev. D* **62**, 084025 (2000).
- [15] T. Gherghetta and A. Pomarol, *Nucl. Phys.* **B586**, 141 (2000); **B602**, 3 (2001).
- [16] N. Arkani-Hamed, T. Gregoire, and J. G. Wacker, *J. High Energy Phys.* **03** (2002) 055; M. A. Luty and R. Sundrum, *Phys. Rev. D* **64**, 065012 (2001); J. Bagger and D. V. Belyaev, *Phys. Rev. D* **67**, 025004 (2003); D. Marti and A. Pomarol, *Phys. Rev. D* **64**, 105025 (2001).
- [17] See, for example, S. J. Huber and Q. Shafi, *Phys. Lett. B* **498**, 256 (2001), and reference therein.
- [18] Z. Lalak, M. Lewicki, M. McGarrie, and P. Olszewski, *J. High Energy Phys.* **11** (2015) 137.
- [19] T. Gherghetta and A. Pomarol, *Nucl. Phys.* **B602**, 3 (2001).
- [20] H. L. Verlinde, *Nucl. Phys.* **B580**, 264 (2000); M. J. Duff, J. T. Liu, and K. S. Stelle, *J. Math. Phys.* **42**, 3027 (2001); M. Cvetič, M. J. Duff, J. T. Liu, H. Lu, C. N. Pope, and K. S. Stelle, *Nucl. Phys.* **B605**, 141 (2001).
- [21] N. Okada and T. Yamada, *Phys. Rev. D* **84**, 035005 (2011).
- [22] N. Maru and N. Okada, *Phys. Rev. D* **70**, 025002 (2004).
- [23] G. F. Giudice and A. Masiero, *Phys. Lett. B* **206**, 480 (1988).
- [24] B. C. Allanach, *Comput. Phys. Commun.* **143**, 305 (2002).
- [25] G. Belanger, F. Boudjema, A. Pukhov, and A. Semenov, *Comput. Phys. Commun.* **149**, 103 (2002); **174**, 577 (2006); **192**, 322 (2015).
- [26] G. Aad *et al.* (ATLAS Collaboration), *J. High Energy Phys.* **10** (2015) 054; CMS Collaboration, Report No. CMS-PAS-SUS-13-019.
- [27] Y. Amhis *et al.* (Heavy Flavor Averaging Group Collaboration), arXiv:1207.1158.
- [28] S. Chatrchyan *et al.* (CMS Collaboration), *Phys. Rev. Lett.* **111**, 101804 (2013).
- [29] D. Asner *et al.* (Heavy Flavor Averaging Group Collaboration), arXiv:1010.1589.
- [30] K. A. Olive *et al.* (Particle Data Group Collaboration), *Chin. Phys. C* **38**, 090001 (2014).
- [31] M. Antonelli *et al.* (FlaviaNet Working Group on Kaon Decays Collaboration), *Nucl. Phys. B, Proc. Suppl.* **181–182**, 83 (2008).
- [32] M. J. Ramsey-Musolf and S. Su, *Phys. Rep.* **456**, 1 (2008); S. Bauman, J. Erler, and M. Ramsey-Musolf, *Phys. Rev. D* **87**, 035012 (2013).
- [33] P. A. R. Ade *et al.* (Planck Collaboration), *Astron. Astrophys.* **594**, A13 (2016).
- [34] D. S. Akerib *et al.* (LUX Collaboration), *Phys. Rev. Lett.* **112**, 091303 (2014); **116**, 161301 (2016).
- [35] M. G. Aartsen *et al.* (IceCube Collaboration), *Phys. Rev. Lett.* **110**, 131302 (2013).